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$D^*\bar{D}^*$ molecule interpretation of $Z_c(4025)$

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Abstract We have used QCD sum rules to study the newly observed charged state $Z_c(4025)$ as a hidden-charm $D^*\bar{D}^*$ molecular state with the quantum numbers $I^G(J^P) = 1^+(1^+)$. Using a $D^*\bar{D}^*$ molecular interpolating current, we have calculated the two-point correlation function and the spectral density up to dimension eight at leading order in α_s. The extracted mass is $m_X = (4.04 \pm 0.24)$ GeV. This result is compatible with the observed mass of $Z_c(4025)$ within the errors, which implies a possible molecule interpretation of this new resonance. We also predict the mass of the corresponding hidden-bottom $B^*\bar{B}^*$ molecular state: $m_{Z_b} = (9.98 \pm 0.21)$ GeV.

1 Introduction

After the observation of charged charmonium-like resonance $Z_c(3900)$ [1], the BESIII Collaboration recently discovered another charged structure $Z_c(4025)$ in the process $e^+e^- \to (D^*\bar{D}^*)^{\pm}\pi^{\mp}$ [2]. This new resonance, which has a mass of $M = (4026.3 \pm 2.6 \pm 3.7)$ MeV, lies very close to the $D^*\bar{D}^*$ threshold. Its width is $\Gamma = (24.8 \pm 5.6 \pm 5.7)$ MeV [2]. To date, the experiment has not determined the quantum numbers of the $Z_c(4025)$ resonance. Since it was observed in both the $D^*\bar{D}^*$ and the $h_c\pi$ channels, the quantum numbers of the charged $Z_c(4025)$ was argued to be $I^G(J^P) = 1^+(1^+)$ while its neutral partner carries negative C-parity [3].

Similar to the other charged charmonium-like states $Z^+(4050)$, $Z^+(4250)$ [4], $Z^+(4430)$ [5] and $Z_c(3900)$ [1], $Z_c(4025)$ cannot be a conventional $c\bar{c}$ state due to the charge it carries. Molecular and tetraquark configurations have recently been used to explore its underlying structure [3]. In Ref. [3], the authors have studied the mass spectrum of $Z_c(4025)$ and its pionic and radiative decays as a $D^*\bar{D}^*$ molecular state using the one-boson-exchange (OBE) model. They have also studied the $Y(4260) \to (D^*\bar{D}^*)^-\pi^+$ decay through the initial-single-pion-emission mechanism in Ref. [6]. $Z_c(4025)$ was also studied as a $[cu][\bar{c}\bar{d}]$ tetraquark with the quantum numbers $J^P = 2^+$ using QCD sum rules [7]. In Ref. [8], the $D^*\bar{D}^*$ $J^P = 1^+$ molecular current with a derivative has been studied in QCD sum rules and the extracted mass coincides with $Z_c(4025)$.

There also exist other theoretical predictions of this new charged structure before its observation by BESIII [9–11]. Ref. [11] studied the $I^G(J^{PC}) = 1^+(1^{+-})$ charmonium-like tetraquark states, and predicted masses near the $D^*\bar{D}^*$ threshold and the possible decay patterns including the open-charm modes $D\bar{D}^*$, $D^*\bar{D}^*$ and other hidden-charm modes. Up to now, theBESIII Collaboration has not reported the $D\bar{D}^*$ decay mode of $Z_c(4025)$. Right now, it seems that the $D^*\bar{D}^*$ molecule interpretation is slightly more natural.

At the hadronic level, the molecular states are commonly assumed to be bound states of two hadrons formed by the exchange of the color-singlet mesons. This configuration is very different from that of the tetraquark states, which are generally bound by the QCD force at the quark–gluon level. In this work, we study $Z_c(4025)$ as a $D^*\bar{D}^*$ molecular state using QCD sum rules approach [12–14].

Within the framework of the QCD sum rule, all the procedures such as the operator product expansion, the calculation of the Wilson coefficient and the Borel transform are very similar for the molecular-type current and tetraquark-type current. In principle, if we exhaust all the possible molecular-type of currents and all the possible tetraquark-type currents, we can rigorously show that these two sets of interpolating currents are equivalent by using a Fierz rearrangement [15,16].

However, there exists an important difference between one single molecular-type current and one single tetraquark-type current. By Fierz rearrangement, every single tetraquark-type current can be expressed as a linear combination of

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several (sometimes up to five) independent molecular-type currents. We can decompose the tetraquark interpolating current into these explicit molecular-type operators. In the single molecular-type QSR, the color flow of the correlation function is quite simple and forms two closed loops. In the tetraquark correlator, there exist additional contributions from the non-diagonal correlator besides the many diagonal correlators as in the molecular-type QSR. Now the color flow is complicated, which is the interference and transition between different molecular structures [17]. In this respect, one well-known example is the light scalar–isoscalar sigma meson. The tetraquark-type current (or their combination/mixing) leads to a better mass prediction than the simple pion–pion molecular current [15].

The paper is organized as follows. In Sect. 2, we calculate the correlation function and spectral density using the $D^*\bar{D}^*$ molecule current. In Sect. 3, we perform a numerical analysis and extract the mass of $Z_c(4025)$. The last section is a brief summary.

2 QCD sum rule and spectral density

The starting point of QCD sum rules is the two-point correlation function

$$\Pi_{\mu\nu}(q^2) = i\int d^4x e^{iq\cdot x}\langle 0|T[J_\mu(x)J_\nu^\dagger(0)]|0\rangle, \quad (1)$$

where $J_\mu(x)$ is the $D^*\bar{D}^*$ molecular interpolating current with $I^G(J^P)=1^+(1^+)$

$$J_\mu = (\bar{q}_a\gamma^\alpha c_a)(\bar{c}_b\sigma_{\alpha\mu}\gamma_5 q_b) - (\bar{q}_a\sigma_{\alpha\mu}\gamma_5 c_a)(\bar{c}_b\gamma^\alpha q_b), \quad (2)$$

in which a, b are color indices and q denotes an up or down quark. In principle, the anti-symmetric tensor operator $\bar{q}_a\sigma_{\alpha\mu}\gamma_5 c_a$ can couple to both $J^P=1^+$ ($\bar{q}_a\sigma_{0i}\gamma_5 c_a$ components) and $J^P=1^-$ ($\bar{q}_a\sigma_{ij}\gamma_5 c_a$ components) channels. However, we can pick out the 1^- piece by multiplication with the vector operator $\bar{c}_b\gamma^\alpha q_b$ so that the molecular operator $(\bar{q}_a\sigma_{\alpha\mu}\gamma_5 c_a)(\bar{c}_b\gamma^\alpha q_b)$ carries the quantum numbers $J^P=1^+$ after contracting the Lorentz index. The molecule current in Eq. (2) contains both the charged components with $(\bar{u}c)(\bar{c}d)$ and $(\bar{d}c)(\bar{c}u)$ pieces and the neutral component with $(\bar{u}c)(\bar{c}u)$ and $(\bar{d}c)(\bar{c}d)$ pieces. For the neutral component, it carries negative C-parity and the quantum numbers should be $I^G(J^{PC})=1^+(1^{+-})$. However, we do not differentiate between u and d quarks in our analysis, so the charged component and the neutral component are the same in QCD sum rules due to isospin symmetry.

The correlation function in Eq. (1) can be written as two independent Lorentz structures since J_μ is not a conserved current:

$$\Pi_{\mu\nu}(q^2) = \left(\frac{q_\mu q_\nu}{q^2} - g_{\mu\nu}\right)\Pi_1(q^2) + \frac{q_\mu q_\nu}{q^2}\Pi_0(q^2), \quad (3)$$

in which the invariant functions $\Pi_1(q^2)$ and $\Pi_0(q^2)$ are related to the spin-1 and spin-0 mesons, respectively. We focus on $\Pi_1(q^2)$ to study the 1^+ channel in this work.

The correlation function in Eq. (1) can be obtained at both the hadron level and the quark–gluon level. To determine the correlation function at the hadron level, we use the dispersion relation

$$\Pi(q^2) = (q^2)^N \int_{4m_c^2}^{\infty} \frac{\rho(s)}{s^N(s-q^2-i\epsilon)}ds + \sum_{n=0}^{N-1} b_n(q^2)^n, \quad (4)$$

where b_n is the unknown subtraction constant which can be removed by taking the Borel transform. The lower limit of integration is the square of the sum of the masses of all current quarks (omitting the light quark mass). $\rho(s)$ is the spectral function

$$\begin{aligned}\rho(s) &\equiv \sum_n \delta(s-m_n^2)\langle 0|J_\mu|n\rangle\langle n|J_\nu^\dagger|0\rangle \\ &= f_X^2\delta(s-m_X^2) + \text{continuum}. \end{aligned} \quad (5)$$

Here we adopt the pole plus continuum parametrization of the hadronic spectral density. The intermediate states $|n\rangle$ must have the same quantum numbers as the interpolating currents J_μ. $|X\rangle$ is the lowest lying resonance with mass m_X and it couples to the current J_μ via the coupling parameter f_X,

$$\langle 0|J_\mu|X\rangle = f_X\epsilon_\mu, \quad (6)$$

where ϵ_μ is the polarization vector ($\epsilon\cdot q=0$).

At the quark–gluon level, the correlation function can be calculated in terms of quark and gluon fields via the operator product expansion (OPE) method. We evaluate the correlation function up to dimension-eight condensate contributions at leading order in α_s using the same technique as in Refs. [11,18–20]. The spectral density is then obtained: $\rho(s)=\frac{1}{\pi}\mathrm{Im}\Pi(q^2)$.

Sum rules for the hadron parameters are established by equating the correlation functions obtained at both the hadron level and the quark–gluon level via quark–hadron duality. The Borel transform is applied to the correlation functions at both levels to remove the unknown constants in Eq. (4), suppress the continuum contribution, and improve the convergence of the OPE series. Using the spectral function defined in Eq. (5), the sum rules can be written as

$$\begin{aligned} f_X^2 m_X^{2k} e^{-m_X^2/M_B^2} &= \int_{4m_c^2}^{s_0} ds e^{-s/M_B^2}\rho(s)s^k \\ &= \mathcal{L}_k\left(s_0, M_B^2\right), \end{aligned} \quad (7)$$

where s_0 is the continuum threshold parameter and M_B is the Borel mass. Then m_X can be extracted by the ratio

$$m_X = \sqrt{\frac{\mathcal{L}_1\left(s_0, M_B^2\right)}{\mathcal{L}_0\left(s_0, M_B^2\right)}}. \tag{8}$$

In the following, we study the lowest lying hadron mass m_X in Eq. (8) as function of the continuum threshold s_0 and Borel mass M_B. We calculate the spectral density at the quark–gluon level including the perturbative term, quark condensate $\langle\bar{q}q\rangle$, gluon condensate $\langle g_s^2 GG\rangle$, quark–gluon mixed condensate $\langle\bar{q}g_s\sigma\cdot Gq\rangle$, four quark condensate $\langle\bar{q}q\rangle^2$, and the dimension-eight condensate $\langle\bar{q}q\rangle\langle\bar{q}g_s\sigma\cdot Gq\rangle$:

$$\rho(s) = \rho^{\text{pert}}(s) + \rho^{\langle\bar{q}q\rangle}(s) + \rho^{\langle GG\rangle}(s) + \rho^{\langle\bar{q}q\rangle^2}(s) + \rho^{\langle\bar{q}Gq\rangle}(s) + \rho^{\langle\bar{q}q\rangle\langle\bar{q}Gq\rangle}(s), \tag{9}$$

where

$$\begin{aligned}
\rho^{\text{pert}}(s) &= \int_{\alpha_{\min}}^{\alpha_{\max}} \mathrm{d}\alpha \int_{\beta_{\min}}^{\beta_{\max}} \mathrm{d}\beta \left[(\alpha+\beta)m_c^2 - \alpha\beta s\right]^3 \\
&\quad\times(1-\alpha-\beta)\left\{\frac{m_c^2(\alpha+\beta-1)(5+\alpha+\beta)}{512\pi^6\alpha^3\beta^3}\right. \\
&\quad\left.+\frac{9(1+\alpha+\beta)\left[(\alpha+\beta)m_c^2-\alpha\beta s\right]}{2048\pi^6\alpha^3\beta^3}\right\}, \\
\rho^{\langle\bar{q}q\rangle}(s) &= -\frac{9m_c\langle\bar{q}q\rangle}{64\pi^4}\int_{\alpha_{\min}}^{\alpha_{\max}} \mathrm{d}\alpha \int_{\beta_{\min}}^{\beta_{\max}} \mathrm{d}\beta(1-\alpha-\beta) \\
&\quad\times\frac{\left[(\alpha+\beta)m_c^2-\alpha\beta s\right]\left[3m_c^2(\alpha+\beta)-7\alpha\beta s\right]}{\alpha\beta^2}, \\
\rho^{\langle GG\rangle}(s) &= \frac{\langle g_s^2 GG\rangle}{1024\pi^6}\int_{\alpha_{\min}}^{\alpha_{\max}} \mathrm{d}\alpha \int_{\beta_{\min}}^{\beta_{\max}} \mathrm{d}\beta(1-\alpha-\beta) \\
&\quad\times\left\{\frac{\left[(\alpha+\beta)m_c^2-\alpha\beta s\right]\left[m_c^2(3+\alpha+\beta)+2\alpha\beta s\right]}{\alpha^2\beta}\right. \\
&\quad+m_c^2(1-\alpha-\beta)\left[\frac{3\left[m_c^2(\alpha+\beta)-2\alpha\beta s\right]}{\alpha^3}\right. \\
&\quad-\frac{(5+\alpha+\beta)\left[m_c^2(4\alpha+3\beta)-3\alpha\beta s\right]}{6\alpha\beta^3} \\
&\quad\left.\left.-\frac{(5+\alpha+\beta)\left[m_c^2(3\alpha+4\beta)-3\alpha\beta s\right]}{6\alpha^3\beta}\right]\right\}, \\
\rho^{\langle\bar{q}Gq\rangle}(s) &= \frac{m_c\langle\bar{q}g_s\sigma\cdot Gq\rangle}{64\pi^4}\int_{\alpha_{\min}}^{\alpha_{\max}} \mathrm{d}\alpha \int_{\beta_{\min}}^{\beta_{\max}} \mathrm{d}\beta \\
&\quad\times\left\{\frac{(1-\alpha-\beta)\left[3m_c^2(\alpha+\beta)-4\alpha\beta s\right]}{\beta^2}\right. \\
&\quad\left.+\frac{(2+7\alpha-2\beta)\left[3m_c^2(\alpha+\beta)-5\alpha\beta s\right]}{2\alpha\beta}\right\}, \\
\rho^{\langle\bar{q}q\rangle^2}(s) &= \frac{5(s+2m_c^2)\langle\bar{q}q\rangle^2}{48\pi^2}\sqrt{1-4m_c^2/s} \qquad (10) \\
\rho^{\langle\bar{q}q\rangle\langle\bar{q}Gq\rangle}(s) &= \frac{\langle\bar{q}q\rangle\langle\bar{q}g_s\sigma\cdot Gq\rangle}{48\pi^2}\int_0^1 \mathrm{d}\alpha \\
&\quad\times\left\{\frac{3m_c^4(3-\alpha)}{\alpha^2(1-\alpha)}\delta'\left[s-\frac{m_c^2}{\alpha(1-\alpha)}\right]\right. \\
&\quad+\frac{m_c^2(3\alpha^3-4\alpha^2-3\alpha+6)}{\alpha(1-\alpha)^2}\delta\left[s-\frac{m_c^2}{\alpha(1-\alpha)}\right] \\
&\quad\left.+(3+2\alpha)H\left[s-\frac{m_c^2}{\alpha(1-\alpha)}\right]\right\},
\end{aligned}$$

in which $\alpha_{\min} = \frac{1-\sqrt{1-4m_c^2/s}}{2}$, $\alpha_{\max} = \frac{1+\sqrt{1-4m_c^2/s}}{2}$, $\beta_{\min} = \frac{\alpha m_c^2}{\alpha s - m_c^2}$, $\beta_{\max} = 1-\alpha$, m_c is the charm quark mass, and $H(\alpha)$ is the Heaviside step function. As is evident from the above expressions, our calculations are of leading order in α_s. Both the quark condensate $\langle\bar{q}q\rangle$ and the quark–gluon mixed condensate $\langle\bar{q}g_s\sigma\cdot Gq\rangle$ are proportional to the charm quark mass m_c. They give important power corrections to the correlation functions. We ignore the chirally suppressed terms proportional to the light quark mass. Based on Ref. [18] the contribution of the three gluon condensate $g_s^3\langle fGGG\rangle$ is expected to be numerically small and has not been included in this work. The dimension-eight condensate $\langle\bar{q}q\rangle\langle\bar{q}g_s\sigma\cdot Gq\rangle$ contains the delta function $\delta\left[s-\frac{m_c^2}{\alpha(1-\alpha)}\right]$ and its derivative. These terms compensate for the singular behavior of the spectral densities at the $s = 4m_c^2$ threshold.

3 Numerical analysis

The following QCD parameters are used in our analysis [21–25]:

$$\begin{aligned}
&m_c(m_c) = (1.23 \pm 0.09)\ \text{GeV}, \\
&m_b(m_b) = (4.20 \pm 0.07)\ \text{GeV}, \\
&\langle\bar{q}q\rangle = -(0.23 \pm 0.03)^3\ \text{GeV}^3, \\
&\langle\bar{q}g_s\sigma\cdot Gq\rangle = -M_0^2\langle\bar{q}q\rangle, \\
&M_0^2 = (0.8 \pm 0.2)\ \text{GeV}^2, \\
&\langle g_s^2 GG\rangle = (0.88 \pm 0.14)\ \text{GeV}^4,
\end{aligned} \tag{11}$$

where the charm and bottom quark masses are the running mass in the $\overline{MS}$ scheme. As mentioned earlier, we set the light quark masses $m_q = 0$ in the analysis. The convention for the mixed condensate is consistent with Refs. [11,18–20], which have a sign difference from some other QCD sum rule studies because of the definition of the coupling constant g_s.

We define the pole contribution (PC) using the sum rules established in Eq. (7),

$$\text{PC}(s_0, M_B^2) = \frac{\mathcal{L}_0\left(s_0, M_B^2\right)}{\mathcal{L}_0\left(\infty, M_B^2\right)}, \tag{12}$$

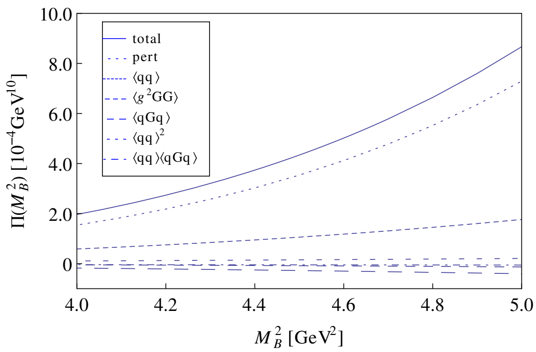

Fig. 1 The convergence of the OPE series with $s_0 \to \infty$

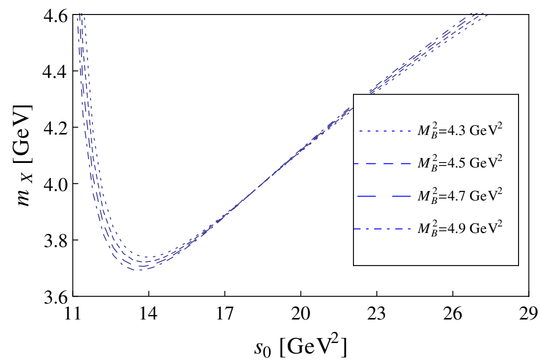

Fig. 2 The variation of m_X with the continuum threshold s_0 in the Borel window $4.3\,\text{GeV}^2 \le M_B^2 \le 4.9\,\text{GeV}^2$

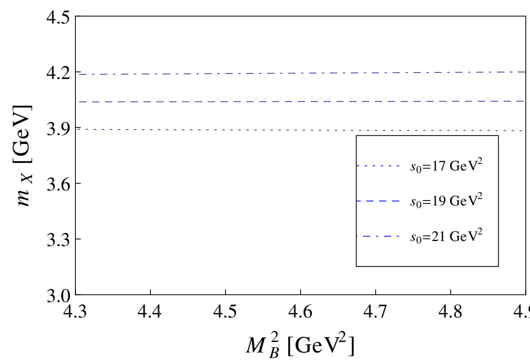

Fig. 3 The variation of m_X with the Borel mass M_B^2 while $s_0 = 17$, 19 and 21 GeV^2

which is the function of the continuum threshold s_0 and the Borel mass M_B. PC represents the lowest lying resonance contribution to the correlation function, which also includes the continuum and higher state contributions with $s_0 \to \infty$.

We begin with the analysis by determining the Borel window. A good mass sum rule requires a suitable working region of the Borel scale M_B. To obtain the lower bound on M_B^2, we let $s_0 \to \infty$ and then study the OPE convergence in Fig. 1. One notes that the quark condensate $\langle\bar{q}q\rangle$ contribution is much bigger than other condensates and is therefore the dominant power correction. Besides the quark condensate, the quark–gluon mixed condensate $\langle\bar{q}g_s\sigma\cdot Gq\rangle$ also gives a significant contribution to the correlation function. From the expression for the spectral density in Eq. (10), the quark condensate and quark–gluon mixed condensate are proportional to the charm quark mass. The gluon condensate $\langle g_s^2 GG\rangle$, four quark condensate $\langle\bar{q}q\rangle^2$, and dimension-eight condensate $\langle\bar{q}q\rangle\langle\bar{q}g_s\sigma\cdot Gq\rangle$ are smaller. However, they also give important corrections to the correlation function and stabilize the mass sum rules. Requiring the quark condensate contribution be less than one third of the perturbative term contribution, while the quark–gluon mixed condensate contribution be less than one third of the quark condensate contribution, we obtain the lower bound on the Borel window $M^2_{\rm min} = 4.3\,\text{GeV}^2$. One may notice from Fig. 1 that the power corrections are small enough in the parameter region $M_B^2 \ge 4.3\,\text{GeV}^2$ so that the OPE convergence is very good.

The continuum threshold s_0 is also an important parameter in QCD sum rules. An optimized choice of s_0 is the value minimizing the variation of the extracted hadron mass m_X with the Borel mass M_B^2. This is achieved by studying the variation of m_X with s_0 in Fig. 2 by varying the value of Borel mass from its lower bound $M^2_{\rm min}$. One notes that these curves with a different value of M_B^2 intersect at $s_0 = 19\,\text{GeV}^2$, around which the variation of m_X with M_B^2 is minimum. Then the upper bound on the Borel mass can be determined by studying the pole contribution defined in Eq. (12). We require that the pole contribution be larger than 10 %, which results in the upper bound on the Borel mass $M^2_{\rm max} = 4.9\,\text{GeV}^2$. We obtain the Borel window $4.3\,\text{GeV}^2 \le M_B^2 \le 4.9\,\text{GeV}^2$ with the threshold value $s_0 = 19\,\text{GeV}^2$.

Now we can perform the QCD sum rule analysis in the Borel window $4.3\,\text{GeV}^2 \le M_B^2 \le 4.9\,\text{GeV}^2$. In Fig. 3, we show the variation of the extracted mass m_X with the Borel mass M_B^2 using continuum thresholds $s_0 = 17$, 19 and 21 GeV^2, respectively. The mass curves are very stable in the Borel window around these threshold values. Finally, we extract the hadron mass:

$$m_X = (4.04 \pm 0.24)\,\text{GeV}, \tag{13}$$

which is very well compatible with the mass of $Z_c(4025)$. This implies the possible $D^*\bar{D}^*$ molecule interpretation of this new resonance.

Using this value of the hadron mass, we also calculate the coupling parameter defined in Eq. (6),

$$f_X = (0.012 \pm 0.005)\,\text{GeV}^5. \tag{14}$$

This parameter represents the strength of the coupling of the current J_μ in Eq. (2) to the $Z_c(4025)$ resonance. The errors of our numerical results in Eqs. (13) and (14) involve the uncertainties in the heavy quark masses and the values of the quark condensate, quark–gluon condensate, and gluon condensate in Eq. (11). Other possible error sources such as

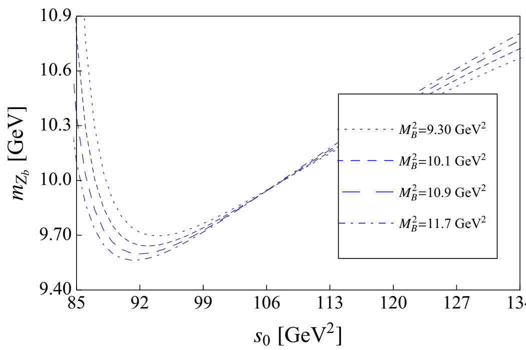

Fig. 4 The variation of m_{Z_b} with the continuum threshold s_0 in the Borel window $9.3\,\text{GeV}^2 \le M_B^2 \le 11.6\,\text{GeV}^2$

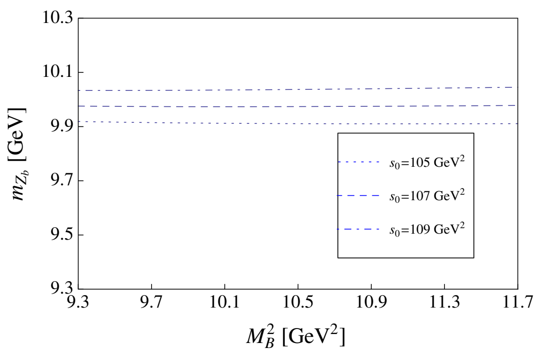

Fig. 5 The variation of m_{Z_b} with the Borel mass M_B^2 while $s_0 = 105$, 107 and $109\,\text{GeV}^2$

truncation of the OPE series, the uncertainty of the continuum threshold value s_0, and the variation of the Borel mass M_B are not taken into account.

We can extend the analysis to the hidden-bottom Z_b system, where Z_b represents a $B^*\bar{B}^*$ molecular state with $I^G(J^P) = 1^+(1^+)$. Using the same interpolating current in Eq. (2), we repeat all the above analysis procedures with the replacement $m_c \to m_b$. To find a suitable working region of the Borel scale, we use the same criteria as in the $D^*\bar{D}^*$ system to study the OPE convergence and pole contribution. We find a Borel window $9.3\,\text{GeV}^2 \le M_B^2 \le 11.6\,\text{GeV}^2$ for the continuum threshold value $s_0 = 107\ \text{GeV}^2$.

We show the Borel curves of the extracted Z_b mass with s_0 and M_B^2 in Figs. 4 and 5, respectively. In Fig. 4, the optimized value of the continuum threshold is chosen as $s_0 = 107\,\text{GeV}^2$, which minimize the variation of the Z_b mass m_{Z_b} with the Borel parameter M_B^2. This result is also shown in Fig. 5, in which the mass curve is very stable as a function of M_B^2 in the obtained Borel window. Considering the same error sources as the $D^*\bar{D}^*$ system, we predict the mass and the coupling parameter of the Z_b state to be

$$m_{Z_b} = (9.98 \pm 0.21)\,\text{GeV}, \tag{15}$$

$$f_{Z_b} = (0.003 \pm 0.001)\,\text{GeV}^5. \tag{16}$$

4 Summary

The BESIII Collaboration has discovered $Z_c(4025)$ in the process $e^+e^- \to (D^*\bar{D}^*)^\pm\pi^\mp$ near the $D^*\bar{D}^*$ threshold. This new structure is a charged resonance and thus cannot be a conventional charmonium state. It is thus a candidate for an exotic hadron state.

In Ref. [8], a $D^*\bar{D}^*$ molecular interpolating current with a derivative operator has been used to investigate the structure of $Z_c(4025)$ in QCD sum rules. In this paper, we use a different hidden-charm $D^*\bar{D}^*$ current with the quantum numbers $I^G(J^P) = 1^+(1^+)$. We have calculated the correlation function and the spectral density up to dimension eight at leading order in α_s, including the perturbative term, quark condensate $\langle\bar{q}q\rangle$, quark–gluon mixed condensate $\langle\bar{q}g_s\sigma\cdot Gq\rangle$, gluon condensate $\langle g_s^2GG\rangle$ and the dimension-eight condensate $\langle\bar{q}q\rangle\langle\bar{q}g_s\sigma\cdot Gq\rangle$ contributions. The quark condensate and the quark–gluon mixed condensate are proportional to the charm quark mass and are a larger contribution than the other condensates. The quark condensate is the dominant power correction to the correlation function. Other condensates are also important because they can improve the OPE convergence and stabilize the mass sum rules.

After performing the QCD sum rule analysis, we extract the hadron mass $m_X = (4.04 \pm 0.24)\,\text{GeV}$ consistent with BESIII's result of the mass of $Z_c(4025)$. Our result supports the $Z_c(4025)$ resonance as an axial-vector $D^*\bar{D}^*$ molecular state. In principle, our result also contains the neutral partner of $Z_c(4025)$ with the quantum numbers $J^{PC} = 1^{+-}$. However, it has the same mass with the charged state in QCD sum rules due to isospin symmetry. We have also studied the corresponding hidden-bottom $B^*\bar{B}^*$ molecular state and predicted the mass $m_{Z_b} = (9.98 \pm 0.21)\,\text{GeV}$. Hopefully our investigation will be useful for the understanding of the structure of the newly observed charged state $Z_c(4025)$ and the future search of its neutral partner.

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References

1. M. Ablikim et al., BESIII collaboration. Phys. Rev. Lett. **110**, 252001 (2013). arXiv:1303.5949 [hep-ex]
2. M. Ablikim et al. (BESIII collaboration) (2013). arXiv:1308.2760 [hep-ex]

3. J. He, X. Liu, Z.-F. Sun, S.-L. Zhu, Eur. Phys. J. C **73**, 2635 (2013). arXiv:1308.2999 [hep-ph]
4. R. Mizuk et al., Belle. Phys. Rev. D **78**, 072004 (2008). arXiv:0806.4098 [hep-ex]
5. S.K. Choi et al., BELLE. Phys. Rev. Lett. **100**, 142001 (2008). arXiv:0708.1790 [hep-ex]
6. X. Wang, Y. Sun, D.-Y. Chen, X. Liu, T. Matsuki (2013). arXiv:1308.3158 [hep-ph]
7. C.-F. Qiao, L. Tang (2013). arXiv:1308.3439 [hep-ph]
8. C.-Y. Cui, Y.-L. Liu, M.-Q. Huang (2013). arXiv:1308.3625 [hep-ph]
9. Z.-F. Sun, J. He, X. Liu, Z.-G. Luo, S.-L. Zhu, Phys. Rev. D **84**, 054002 (2011). arXiv:1106.2968 [hep-ph]
10. D.-Y. Chen, X. Liu, Phys. Rev. D **84**, 034032 (2011). arXiv:1106.5290 [hep-ph]
11. W. Chen, S.-L. Zhu, Phys. Rev. D **83**, 034010 (2011). arXiv:1010.3397 [hep-ph]
12. M.A. Shifman, A.I. Vainshtein, V.I. Zakharov, Nucl. Phys. B **147**, 385 (1979)
13. L.J. Reinders, H. Rubinstein, S. Yazaki, Phys. Rep. **127**, 1 (1985)
14. K.A. Colangelo, Pietro, Front. Particle Phys. **3** (2000). arXiv:hep-ph/0010175
15. H.-X. Chen, A. Hosaka, S.-L. Zhu, Phys. Rev. D **76**, 094025 (2007). arXiv:0707.4586 [hep-ph]
16. C.-K. Jiao, W. Chen, H.-X. Chen, S.-L. Zhu, Phys. Rev. D **79**, 114034 (2009)
17. M. Nielsen, F.S. Navarra, S.H. Lee, Phys. Rep. **497**, 41 (2010). arXiv:0911.1958 [hep-ph]
18. W. Chen, S.-L. Zhu, Phys. Rev. D **81**, 105018 (2010). arXiv:1003.3721 [hep-ph]
19. M.-L. Du, W. Chen, X.-L. Chen, S.-L. Zhu, Chin. Phys. C **37**, 033104 (2013). arXiv:1203.5199 [hep-ph]
20. M.-L. Du, W. Chen, X.-L. Chen, S.-L. Zhu, Phys. Rev. D **87**, 014003 (2013). arXiv:1209.5134 [hep-ph]
21. J. Beringer et al., Particle data group. Phys. Rev. D **86**, 010001 (2012)
22. M. Eidemuller, M. Jamin, Phys. Lett. B **498**, 203 (2001). arXiv:hep-ph/0010334
23. M. Jamin, A. Pich, Nucl. Phys. Proc. Suppl. **74**, 300 (1999). arXiv:hep-ph/9810259
24. M. Jamin, J.A. Oller, A. Pich, Eur. Phys. J. C **24**, 237 (2002). arXiv:hep-ph/0110194
25. A. Khodjamirian, Th. Mannel, N. Offen, Y.- M. Wang, Phys. Rev. D **83**, 094031 (2011)